

LHC Accessible Second Higgs Boson in the Left-Right Model

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Abstract

A second Higgs doublet arises naturally as a parity partner of the standard model (SM) Higgs, once SM is extended to its left-right symmetric version (LRSM) to understand the origin of parity violation in weak interactions as well as to accommodate small neutrino masses via the seesaw mechanism. The flavor changing neutral Higgs (FCNH) effects in the minimal version of this model (LRSM), however, push the second Higgs mass to more than 15 TeV making it inaccessible at the LHC. Furthermore since the second Higgs mass is directly linked to the W_R mass, discovery of a “low” mass W_R ($M_{W_R} \leq 5 - 6$ TeV) at the LHC would require values for some Higgs self couplings larger than one. In this paper we present an extension of LRSM by adding a vector-like $SU(2)_R$ quark doublet which weakens the FCNH constraints allowing the second Higgs mass to be near or below TeV and a third neutral Higgs below 3 TeV for a W_R mass below 5 TeV. It is then possible to search for these heavier Higgs bosons at the LHC, without conflicting with FCNH constraints. A right handed W_R mass in the few TeV range is quite natural in this class of models without having to resort to large scalar coupling parameters. The CKM mixings are intimately linked to the vector-like quark mixings with the known quarks, which is the main reason why the constraints on the second Higgs mass is relaxed. We present a detailed theoretical and phenomenological analysis of this extended LR model and point out some tests as well as its potential for discovery of a second Higgs at the LHC. Two additional features of the model are: (i) a 5/3 charged quark and (ii) a fermionic top partner with masses in the TeV range.

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I. INTRODUCTION

With the discovery of the standard model Higgs boson at the Large Hadron Collider [1], attention has now shifted in part to the search for a second heavier Higgs boson [2, 3]. While the 125 GeV Higgs boson has confirmed the standard model, the second Higgs field is likely to provide strong clues to the nature of new physics beyond the standard model. For instance, a second Higgs boson is a natural part of several extensions of SM e.g. minimal supersymmetric standard model (MSSM), Peccei-Quinn extension to solve the strong CP problem as well as models with spontaneous CP violation as in multi-Higgs extensions of SM[4]. Another class of models where also a second Higgs doublet is forced on us by gauge symmetry is the left-right symmetric (LRSM) extension of SM [5], which provides a way[6] to understand the small neutrino masses via the seesaw mechanism. The second Higgs doublet in this model is the parity partner of the SM Higgs and is dictated by the gauge group of the LRSM, $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The minimal version of this model (MLRSM) is defined as the one with parity symmetric fermion and Higgs assignments with scalar bidoublet field $\phi(2, 2, 0)$ giving masses to charged fermions and the $\Delta_L(3, 1, 2) \oplus \Delta_R(1, 3, 2)$ breaking the $SU(2)_R \times U(1)_{B-L}$ symmetry as well as implementing the seesaw mechanism. In this minimal version, the SM Higgs field (ϕ_{SM}) is part of a bi-doublet field ϕ which contains two SM doublets, the second being the parity partner of ϕ_{SM} . It turns out however that in MLRSM, gauge invariance also restricts the coupling of the second Higgs to the quarks in such a way that it leads to large flavor changing neutral Higgs effects unless its mass is more than about 15 TeV [7]. This pushes the second Higgs boson beyond the reach of LHC but perhaps more importantly, if a right handed W_R is discovered at the LHC with mass below 5-6 TeV, some self scalar coupling parameter in the potential must be larger than one, causing some tension.

This also raises the following more practical question: suppose LHC discovers a second Higgs boson with a few TeV (or less) mass; in that case, should the search for W_R boson at the LHC [8] stop? The minimal LRMS would say “yes” since the second TeV-ish Higgs boson would rule this model out. What we point out in this paper is that a second TeV-ish neutral Higgs boson does not necessarily rule out the general class of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ extensions of SM since as we show in this paper, there exist non-minimal left-right models with few TeV mass W_R which can accommodate a near TeV mass Higgs boson without conflicting with FCNH constraints from meson-anti-meson mixings. The search for TeV scale W_R boson at the 14 TeV LHC should therefore continue even if a TeV-ish Higgs boson is discovered. In fact we point out that in the example we propose, the search for W_R should most likely continue in the tri-lepton mode rather than the $\ell^\pm \ell^\pm jj$ mode [9, 10] currently being used. While this may not be a generic feature of such extended models, it may be something to keep in mind until a different example is found. The goal of this paper is to provide an existence proof by example of such a model¹.

It was pointed few years ago that if we imagine the TeV scale left-right model as an effective theory, then one can add higher dimensional ($d = 6$) operators to the theory that involve additional Higgs fields and the minimal LRSM fermions which can help to bring consistency between the TeV scale Higgs mass with FCNH effects [12]. The example provided in this paper is a UV complete version of the model which requires adding a single vector-like $SU(2)_R$ quark doublet and an extra Higgs bi-doublet carrying $B - L$. We show

¹ For an alternative approach, see [11].

that this model reproduces the CKM mixing in the quark sector with tiny deviations from unitarity and we then analyze the FCNH effects which exist only in the up-quark sector, i.e. in the $D^0 - \bar{D}^0$ mixing. We find that for the second Higgs boson mass near a TeV, we can satisfy the FCNH bounds. There are no FCNH effects in the down quark sector due to the symmetries of the model. As part of our study of phenomenological implications of the model, we discuss production and decay modes of the second heavy Higgs.

The basic outline of our strategy is as follows. The root of the FCNH bound on the second Higgs in the minimal LRSM is that the bi-doublet field ϕ and its conjugate $\tilde{\phi}$ both couple to quark doublets as $\mathcal{L}_Y = h\bar{Q}_L\phi Q_R + \tilde{h}\bar{Q}_L\tilde{\phi}Q_R + h.c..$ When quark mass matrices are diagonalized to generate the V_{CKM} , the neutral component of the second Higgs doublet has off diagonal couplings which then give rise to the Flavor changing effects. The first step to cure this problem is to prevent the $\tilde{\phi}$ Yukawa coupling to quarks, while allowing the ϕ coupling. This however leads to $V_{CKM} = 1$ so that all quark mixings vanish. Our suggestion to cure this problem is to introduce one set of vector like quarks which are such that generate the quark mixings only in the up sector. As a result, the only FCNH effect we have to consider is the $D^0 - \bar{D}^0$ mixing. Since CKM mixings arise due to small mixings with the vector like quarks, the resulting constraint on the second neutral Higgs is much weaker. The presence of the extra bi-doublet also leads to a third Higgs field with mass below 3 TeV for a W_R mass less than 5 TeV. While we do not go into details of the lepton sector of our model, we note that small neutrino masses most likely arise in this model from the inverse seesaw mechanism [13]. This in turn implies that the search for W_R should focus on the tri-lepton mode [14].

The paper is organized as follows: In sec. 2, we present the particle content and the fermion sector of the model and then obtain a parameter range where the correct CKM mixings arise; in sec. 3, we discuss the Higgs potential, its minimization to obtain the neutral Higgs spectrum; in sec. 4, we find the FCNH constraints on the Higgs masses. In sec. 5, we choose some bench mark points of the model and discuss the LHC signal for the Higgs fields of the model. In sec. 6, we briefly touch on the lepton sector of the model, more specifically, the origin of neutrino masses. In sec. 7, we discuss some other phenomenological implications as well as comment on the grand unification prospects for the model. We present a summary of our results in the final section 8. In appendix A, we present an example of quark mixing solution without CP violation and in appendix B, we give details of the potential minimization and neutral scalar mass diagonalization.

II. EXTENDED LEFT-RIGHT MODEL

In addition to the usual left-right fermion doublets $Q_{a,L,R}^T = (u_a, d_a)_{L,R}$ and $\ell_{a,L,R}^T = (\nu_a, e_a)_{L,R}$ with obvious $U(1)_{B-L}$ quantum numbers, we add an $SU(2)_R$ vector like quark doublet $Q'^T \equiv (T, t')$ with $B - L = \frac{7}{3}$. Because of the exotic B-L assignment, the vector-like quark T has electric charge $5/3$ and t' has $Q = 2/3$ like the top quark. We also note that the model at the TeV scale does not respect discrete parity invariance. However, adding extra heavy $SU(2)_L$ vector-like quarks and parity odd Higgs fields, the model can be made parity invariant at a high scale. As a result of this high scale breakdown of parity there is no type II contribution to neutrino masses.

The Higgs sector of the extended LRSM model suggested here consists of the following Higgs fields, with their quantum numbers under the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge

symmetry given within bracket:

$$\begin{aligned}
\phi &= \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} && \in (2, 2, 0), \\
\rho &= \begin{pmatrix} \rho_1^+ & \rho^{++} \\ \rho^0 & \rho_2^+ \end{pmatrix} && \in (2, 2, 2), \\
\Delta_L &= \begin{pmatrix} \delta_L^+/\sqrt{2} & \delta_L^{++} \\ \delta_L^0 & -\delta_L^+/\sqrt{2} \end{pmatrix} && \in (3, 1, 2), \\
\Delta_R &= \begin{pmatrix} \delta_R^+/\sqrt{2} & \delta_R^{++} \\ \delta_R^0 & -\delta_R^+/\sqrt{2} \end{pmatrix} && \in (1, 3, 2).
\end{aligned} \tag{1}$$

Under the $SU(2)_L \times SU(2)_R$ gauge symmetry, these fields transforms as

$$\begin{aligned}
\phi &\rightarrow U_L \phi U_R^\dagger, & \tilde{\phi} &\rightarrow U_L \tilde{\phi} U_R^\dagger, \\
\rho &\rightarrow U_L \rho U_R^\dagger, & \tilde{\rho} &\rightarrow U_L \tilde{\rho} U_R^\dagger, \\
\Delta_L &\rightarrow U_L \Delta_L U_L^\dagger, & \Delta_L^\dagger &\rightarrow U_L \Delta_L^\dagger U_L^\dagger, \\
\Delta_R &\rightarrow U_R \Delta_R U_R^\dagger, & \Delta_R^\dagger &\rightarrow U_R \Delta_R^\dagger U_R^\dagger.
\end{aligned} \tag{2}$$

where $\tilde{\varphi} = -i\sigma_2 \varphi^* i\sigma_2$ ($\varphi = \phi, \rho$), and $U_{L,R}$ are, respectively, the general $SU(2)_L$ and $SU(2)_R$ unitarity transformations.

We assume the theory to be invariant under a discrete Z_4 symmetry so that simultaneous coupling of ϕ and $\tilde{\phi}$ couplings to the SM quarks is forbidden naturally. The complete set of transformations of all the fields in the extended LRSM under the discrete Z_4 symmetry are given below:

$$\begin{aligned}
\phi &\rightarrow i\phi, & Q_L &\rightarrow Q_L, & \ell_L &\rightarrow \ell_L, \\
\rho &\rightarrow i\rho, & Q_R &\rightarrow -iQ_R, & \ell_R &\rightarrow i\ell_R. \\
\Delta_R &\rightarrow -\Delta_R, & Q'_{L,R} &\rightarrow iQ'_{L,R},
\end{aligned} \tag{3}$$

Consequently, for the two scalars ϕ and ρ ,

$$\tilde{\varphi} \rightarrow -i\tilde{\varphi}. \tag{4}$$

The Yukawa couplings of the model read, under the discrete symmetry,

$$\begin{aligned}
-\mathcal{L}_Y &= \bar{Q}_L h_q \phi Q_R + \bar{Q}_L y_f \tilde{\rho} Q'_R + \bar{Q}'_L y_g \Delta_R Q_R + M \bar{Q}'_L Q'_R \\
&\quad + \bar{\ell}_L h_\ell \tilde{\phi} \ell_R + y_R \ell_R \ell_R \Delta_R + \text{h.c.} .
\end{aligned} \tag{5}$$

After the scalars get non-vanishing vevs (for simplicity we assume all the vevs are real),

$$\begin{aligned}
\langle \phi \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} \kappa_1 & 0 \\ 0 & \kappa_2 \end{pmatrix}, \\
\langle \rho \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_\rho & 0 \end{pmatrix}, \\
\langle \Delta_R \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix},
\end{aligned} \tag{6}$$

the quark mass matrices read

$$\begin{aligned}\mathcal{M}_d &= \frac{1}{\sqrt{2}}\kappa_2 h_q, \\ \mathcal{M}_u &= \frac{1}{\sqrt{2}} \begin{pmatrix} h_q \kappa_1 & y_f v_\rho \\ y_g v_R & \sqrt{2}M \end{pmatrix} = \begin{pmatrix} h & f \\ g & M \end{pmatrix},\end{aligned}\tag{7}$$

where h_q is a 3×3 matrix which can be chosen to be diagonal by a choice of quark basis, and consequently

$$\mathcal{M}_d = \text{diag}\{m_d, m_s, m_b\}.\tag{8}$$

Then the quark mixings as well as CP violation, i.e. a nontrivial CKM matrix, comes from the f and g parameters, or more specifically from the Yukawa couplings y_f and y_g . Note that the y_f and y_g are not related by left symmetry since they are coupled to different Higgs bosons not related by parity at high scale. This means that the left and right handed quark mixing angles will be very different from each other.

A. CKM fit in the model

As we see from the previous sub-section, the quark mass matrices are very highly constrained and therefore it is a priori not clear that the model will reproduce the correct quark masses and CKM mixings for reasonable choice of parameters. Below we show that this is indeed the case. The starting point of quark mixing is the 4×4 up-type mass matrix in the extended LRSM,

$$\mathcal{M}_u = \begin{pmatrix} h_1 & 0 & 0 & f_1 \\ 0 & h_2 & 0 & f_2 \\ 0 & 0 & h_3 & f_3 \\ g_1 & g_2 & g_3 & M \end{pmatrix}.\tag{9}$$

This matrix $\mathcal{M}_u \mathcal{M}_u^\dagger$ can be diagonalized by the 4×4 CKM matrix V_{CKM4} , in the basis with diagonal down-type quark mass matrix,

$$V_{\text{CKM4}}(\mathcal{M}_u \mathcal{M}_u^\dagger)V_{\text{CKM4}}^\dagger = \text{diag}\{m_u^2, m_c^2, m_t^2, m_t'^2\},\tag{10}$$

where m_t' is the mass for the introduced sequential 4th up-type quark. Equivalently,

$$\mathcal{M}_u \mathcal{M}_u^\dagger = V_{\text{CKM4}}^\dagger \text{diag}\{m_u^2, m_c^2, m_t^2, m_t'^2\} V_{\text{CKM4}}.\tag{11}$$

On the LHS of the equation, in our extended LRSM,

$$\begin{aligned}h_1 &= r m_d, \\ h_2 &= r m_s, \\ h_3 &= r m_b,\end{aligned}\tag{12}$$

where r is the vev ratio κ_1/κ_2 . As the up-type quark mass matrix is diagonalized at the TeV scale, i.e. the new physics scale of our model, we use the RGE-evolved down-type and

up-type quark masses in our fit [15],

$$\begin{aligned} m_u(\text{TeV}) &= 2.6 \text{ MeV}, & m_d(\text{TeV}) &= 2.5 \text{ MeV}, \\ m_c(\text{TeV}) &= 0.53 \text{ GeV}, & m_s(\text{TeV}) &= 47 \text{ MeV}, \\ m_t(\text{TeV}) &= 150.7 \text{ GeV}, & m_b(\text{TeV}) &= 2.43 \text{ GeV}. \end{aligned} \quad (13)$$

On the RHS of Eq. (11), the unitary of the 3×3 CKM matrix is in good agreement with observations, leaving little room for mixing with heavy quarks, i.e. [16]

$$\begin{aligned} |V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 &= 0.9999 \pm 0.0006, \\ |V_{cd}|^2 + |V_{cs}|^2 + |V_{cb}|^2 &= 1.067 \pm 0.047, \\ |V_{td}|^2 + |V_{ts}|^2 + |V_{tb}|^2 &= 1. \pm 0.000137. \end{aligned} \quad (14)$$

In Appendix A, We give a toy fit of the CKM matrix without any CP violation, which however reveal some specific feature of the fitting, e.g. r is required to be of order 10, and the f_i parameters are generally small while g_i are TeV scale parameters. In the realistic CP violating case, M is kept as a real parameter, while f_j and g_j are required to have non-vanishing phases denoted as

$$\begin{aligned} f_j &\rightarrow f_j e^{i\alpha_j}, \\ g_j &\rightarrow g_j e^{i\beta_j}, \end{aligned} \quad (15)$$

then the LHS of Eq. (11) reads

$$\mathcal{M}_u \mathcal{M}_u^\dagger = \begin{pmatrix} h_1^2 + f_1^2 & f_1 f_2 e^{i(\alpha_1 - \alpha_2)} & f_1 f_3 e^{i(\alpha_1 - \alpha_3)} & f_1 M e^{i\alpha_1} + g_1 h_1 e^{-i\beta_1} \\ f_1 f_2 e^{-i(\alpha_1 - \alpha_2)} & h_2^2 + f_2^2 & f_2 f_3 e^{i(\alpha_2 - \alpha_3)} & f_2 M e^{i\alpha_2} + g_2 h_2 e^{-i\beta_2} \\ f_1 f_3 e^{-i(\alpha_1 - \alpha_3)} & f_2 f_3 e^{-i(\alpha_2 - \alpha_3)} & h_3^2 + f_3^2 & f_3 M e^{i\alpha_3} + g_3 h_3 e^{-i\beta_3} \\ f_1 M e^{-i\alpha_1} + g_1 h_1 e^{i\beta_1} & f_2 M e^{-i\alpha_2} + g_2 h_2 e^{i\beta_2} & f_3 M e^{-i\alpha_3} + g_3 h_3 e^{i\beta_3} & g_1^2 + g_2^2 + g_3^2 + M^2 \end{pmatrix}. \quad (16)$$

One representative solution we find for the CP violating case is the following:

$$\begin{aligned} h_1 &= 0.037, \\ h_2 &= 0.346, \\ h_3 &= 13.1; \end{aligned} \quad (17)$$

for the f parameters (in unit of GeV) and its phases (in unit of radian)

$$\begin{aligned} f_1 &= 1.38, \quad \alpha_1 = -2.80, \\ f_2 &= 6.32, \quad \alpha_2 = -0.0499, \\ f_3 &= 158, \quad \alpha_3 = -3.17; \end{aligned} \quad (18)$$

for the g parameters (in unit of GeV) and its phases (in unit of radian)

$$\begin{aligned} g_1 &= 1450, \quad \beta_1 = 1.64, \\ g_2 &= 2398, \quad \beta_2 = -2.93, \\ g_3 &= 573, \quad \beta_3 = 2.43; \end{aligned} \quad (19)$$

and the last parameter, in unit of GeV,

$$M = 905. \quad (20)$$

For the h parameters, if we choose the mass ratio $r = 5$, then the down-type quark masses we need as input are, respectively,

$$\begin{aligned} m_d(\text{TeV}) &= 7.57 \text{ MeV} , \\ m_s(\text{TeV}) &= 69.4 \text{ MeV} , \\ m_b(\text{TeV}) &= 2.62 \text{ GeV} . \end{aligned} \quad (21)$$

which are consistent with their experimental values. With these parameters we can fit successfully all the up-type quark masses, left-handed quark mixing angles and CP violation phase. Furthermore, much like in the CP conserving case, the mixing between the SM up-type quarks with t' are very small,

$$\begin{aligned} s_{14} &= 0.000025 , \\ s_{24} &= 0.00012 , \\ s_{34} &= 0.0165 . \end{aligned} \quad (22)$$

As a direct consequence of the large g_i parameters, however, contrary to the minimal LRSM case [7], the right-handed quark mixings are generally very large, e.g. in the fit above,

$$\begin{pmatrix} 0.142 + 0.779i & 0.513 + 0.0444i & -0.32744 - 0.0112i & -0.0272 \\ -0.0286 + 0.337i & 0.101 - 0.128i & 0.923 + 0.0314i & 0.0765 \\ -0.0171 + 0.158i & -0.255 - 0.0654i & 0.0375 + 0.0425i & -0.950 \\ -0.0336 + 0.482i & -0.781 - 0.169i & -0.145 + 0.124i & 0.302 \end{pmatrix} . \quad (23)$$

Due to the large W_R mass, however, they do not lead to any conflict with observations.

III. SPECTRUM OF HIGGS FIELDS

Now that the model can reproduce the quark masses and mixings, we move on to discuss the masses and decay properties of the second and other neutral Higgs in the model. Our interest is primarily in the second Higgs which consists predominantly of the components $\text{Re}\phi_{1,2}^0$ and $\text{Re}\rho^0$. Our model has four complex neutral Higgs fields $\phi_{1,2}^0, \rho^0, \delta_R^0$ which will in general mix among themselves. We have to find the mass eigenstates. To proceed with this study, we analyze the gauge and Z_4 invariant Higgs potential Eq. (B1) in Appendix B allowing only for two soft Z_4 breaking terms (Eq. (B2)). We then minimize the Higgs potential to obtain the desired minima given in Eq. (6) above and use them to find the spectrum of neutral Higgs fields, which are mass eigenstates [17]. Note that unlike the minimal LRSM, there is no Δ_L field. We adopt high scale D-parity breaking to push this field to very high scale [18], which also eliminates the type II contributions to neutrino masses that could otherwise be “large” in TeV LR models.

For numerical fit, after setting a sum of two scalar couplings (denoted by γ' , cf. Eq. (B5)) $\gamma' = 0.5$, we get a scalar H_R with mass v_R at the leading order. By assigning appropriate

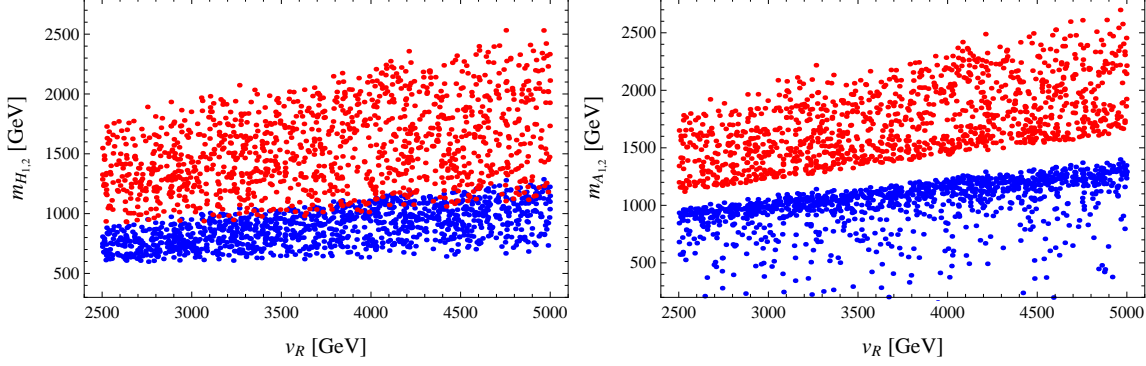


FIG. 1: The masses of $H_{1,2}$ and $A_{1,2}$ as function of v_R . The masses of A_1 and H_1 are in blue and A_2 and H_2 are in red.

values to other relevant coefficients and with the $SU(2)_L$ breaking vevs satisfying the constraint²

$$\kappa_1^2 + \kappa_2^2 + v_\rho^2 = v^2 = (246 \text{ GeV})^2, \quad (24)$$

we can easily get a 125 GeV Higgs recently observed at the LHC; we can also obtain the corresponding masses of the other two CP-even states, shown in Fig. 1 as functions of v_R . It is obvious that $m_{H_1} \sim \sqrt{v v_R}$ and H_2 has a somewhat larger and broader mass range³. In the fit we found that α'_2 is required to be large ~ 10 , as we argued above. It should be stressed, however, that a large α'_2 does not necessarily mean the invalidity of perturbation theory, since it is the linear summation of five independent scalar self couplings, cf. Eq. (B5). Furthermore, to get a 125 GeV Higgs, κ_1 is found to be large, $\sim (130, 210)$ GeV, which means that generally both $\text{Re}\phi_1^0$ and $\text{Re}\rho_0$ contribute substantially to the SM Higgs.

As an explicit example, when we set the vevs

$$\begin{aligned} \kappa_1 &= 183 \text{ GeV}, \\ \kappa_2 &= 36.6 \text{ GeV}, \\ v_\rho &= 160 \text{ GeV}, \\ v_R &= 3305 \text{ GeV}, \end{aligned} \quad (25)$$

and the quartic couplings (all defined in Appendix B)

$$\begin{aligned} \alpha'_2 &= 10.1, \\ \gamma' &= 0.5, \\ y'_1 &= 0.588, \end{aligned}$$

² Numerical analysis reveals that only about one per cent of the SM higgs is from its mixing with the right-handed Higgs δ_R^0 . As a leading order approximation, we can neglect this contribution.

³ In the numerical fit, we set explicitly $\gamma' = 0.5$ and other couplings defined in the Appendix as $\alpha'_2 \lesssim 30$, and $y'_1, z' \lesssim 1$. The vevs κ_1 is taken as a free parameters in the range of $[53, 220]$ GeV (the two extreme values are determined by requiring that the Yukawa coupling in the quark mass matrices are not too large); we choose $\kappa_2 = \kappa_1/5$ (this ratio respects our numerical fit of the 4×4 up-type mass matrix with CP violation) and $v_\rho = \sqrt{v^2 - \kappa_1^2 - \kappa_2^2}$. We select the sets of parameters which predict the lightest Higgs mass $m_h \in [123, 127]$ GeV. For the pseudoscalar mass plot, we set explicitly $M' = v$, and $0 < \alpha_{2,6} \lesssim 3$.

$$z' = 0.064, \quad (26)$$

we can get the SM Higgs (the lightest one) with mass 125 GeV, with help of the rotation matrix connecting the flavor states and the mass eigenstates (the superscript S stands for “scalar”)

$$\begin{pmatrix} h \\ H_1 \\ H_2 \\ H_R \end{pmatrix} = U^S \begin{pmatrix} \text{Re}\phi_1^0 \\ \text{Re}\phi_2^0 \\ \text{Re}\rho^0 \\ \text{Re}\delta_R^0 \end{pmatrix} \quad (27)$$

with values

$$U^S = \begin{pmatrix} 0.762 & 0.086 & 0.640 & -0.029 \\ 0.624 & -0.354 & -0.695 & -0.007 \\ 0.166 & 0.930 & -0.325 & -0.041 \\ 0.0339 & 0.0386 & 0.00065 & 0.998 \end{pmatrix}. \quad (28)$$

(Note that this choice of the $\kappa_{1,2}$ gives the realistic CP violating CKM mixings cf. line after Eq. (20).) In this case, the masses of the three heavier states are, respectively,

$$\begin{aligned} m_{H_1} &= 812 \text{ GeV}, \\ m_{H_2} &= 1335 \text{ GeV}, \\ m_{H_R} &= 3309 \text{ GeV}. \end{aligned} \quad (29)$$

IV. FCNH EFFECTS CONSTRAINT ON THE SECOND HIGGS BOSON MASS

To investigate the FCNH constraints, we look at the couplings of the neutral Higgs fields in the Yukawa coupling given in Eq. (5). For the down-type quarks, all the couplings to the scalars are proportional to the diagonal \mathcal{M}_d and therefore they do not introduce no flavor changing neutral couplings. However, for the up-type quarks, the situation is completely different. The neutral Higgs couplings for the up sector can be written as:

$$-\mathcal{L}_Y^0 = \bar{u}_L(h_q\phi_1^0)u_R + \bar{u}_L(h_f\rho^{0*})t'_R + \bar{t}'_L(h_g\delta_R^0)u_R + M\bar{t}'_Lt'_R. \quad (30)$$

Using the fact that the vevs of these fields contribute to up quark mass matrix \mathcal{M}_u , we can write \mathcal{L}_Y^0 as:

$$-\mathcal{L}_Y^0 = \bar{\mathcal{U}}_L\mathcal{M}_u\mathcal{U}_R + (\sqrt{2}G_F)^{1/2}\bar{\mathcal{U}}_L\mathcal{M}'_uh\mathcal{U}_R + (\sqrt{2}G_F)^{1/2}\bar{\mathcal{U}}_L\mathcal{M}''_uH_1\mathcal{U}_R + \dots, \quad (31)$$

where

$$\mathcal{U} = \begin{pmatrix} u_a \\ t' \end{pmatrix} \quad (32)$$

and

$$\mathcal{M}'_u = \begin{pmatrix} c'_0h & c'_1f \\ c'_2g & 0 \end{pmatrix}, \quad (33)$$

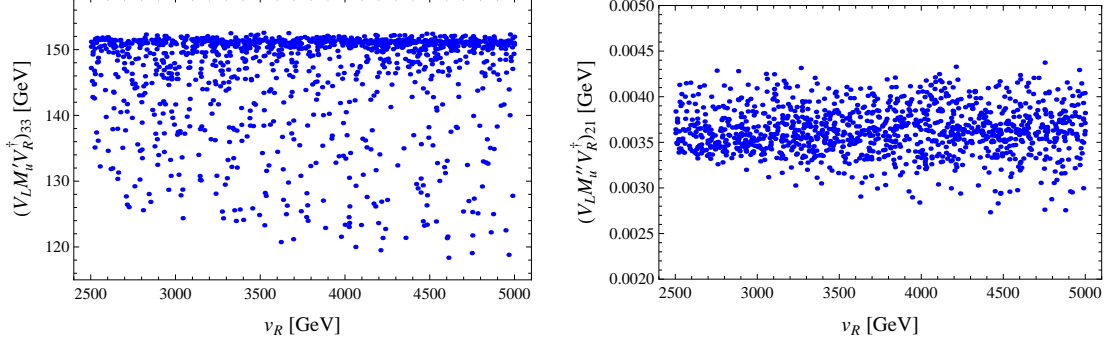


FIG. 2: Distributions of $(V_L M'_u V_R^\dagger)_{33}$ and $(V_L M''_u V_R^\dagger)_{21}$ as functions of v_R , in unit of GeV.

$$\mathcal{M}''_u = \begin{pmatrix} c''_0 h & c''_1 f \\ c''_2 g & 0 \end{pmatrix}. \quad (34)$$

In the numerical example given above,

$$\begin{aligned} c'_0 &= \frac{v}{\kappa_1} U_{11}^S = 1.025, & c''_0 &= \frac{v}{\kappa_1} U_{21}^S = 1.58, \\ c'_1 &= \frac{v}{v_\rho} U_{13}^S = 0.984, & c''_1 &= \frac{v}{v_\rho} U_{23}^S = -0.620, \\ c'_2 &= \frac{v}{v_R} U_{14}^S = -0.0022, & c''_2 &= \frac{v}{v_R} U_{24}^S = -0.00093. \end{aligned} \quad (35)$$

When we transform the up-type quarks into their mass eigenstates, the coupling of SM higgs h couples to the quarks via the matrix

$$V_L \mathcal{M}'_u V_R^\dagger \simeq \begin{pmatrix} 0.0301 & 0.000074 & -0.00098 & -0.0736 \\ -0.000083 & 0.543 & -0.00323 & -0.372 \\ -0.176 & 0.497 & 148. & -49.0 \\ -0.0513 & 0.144 & -4.36 & -5.20 \end{pmatrix}, \quad (36)$$

where V_L and V_R are, respectively, the left- and right-handed quark mixing matrices. Here follow two comments:

- It is transparent from the left panel of Fig. 2 that the coupling to top quark (and the up and charm quarks, given potential corrections to the up quark Yukawa coupling) is mostly about the same as in the SM (note that when evaluated at TeV scale, $m_t \simeq 150$ GeV), and the coupling to the fourth heavy quark is tiny and its contribution to Higgs production at LHC can be safely neglected, with a contribution of $\sim 10^{-3}$ to 10^{-4} .
- It seems that the flavor-changing coupling of the SM Higgs is very small, the largest one of which being the element involving the flavor changing Higgs coupling to the top and charm quark (we do not consider the elements relevant to the heavy fourth quark). The element $(V_L M'_u V_R^\dagger)_{32}$ imply a Yukawa coupling y_{tch} of order $10^{-4} - 10^{-2}$, as shown in Fig. 3. There has been much discussion on constraints of such flavor-changing coupling from top and Higgs experimental data [19–21]. The ATLAS collaboration recently got an upper limit of per cent level on the non-standard top decay $t \rightarrow hc$ channel [22], which imply an upper limit of order 0.1 on the Yukawa coupling y_{tch} . The authors of [23] assert that the processes $pp \rightarrow t\bar{j}h, \bar{t}jh$ at LHC can produce a more stringent constraint, of the order of 10^{-3} . As far as our model goes, it is evident from Fig. 3 that there exists a large region of parameter space which respects the experimental constraints on the flavor changing top couplings.

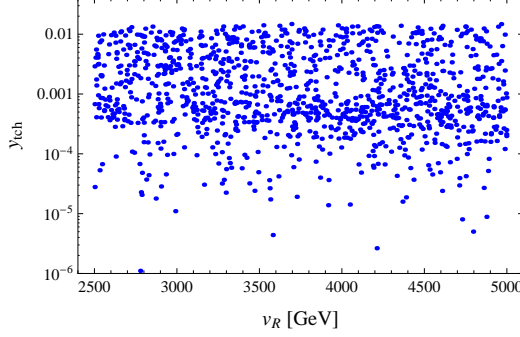


FIG. 3: Flavor changing Yukawa coupling y_{tch} as function of v_R .

In the example, the coupling to the second Higgs H_1 is proportional to

$$V_L \mathcal{M}_u'' V_R^\dagger \simeq \begin{pmatrix} 0.0235 & 0.003 & -0.0408 & -0.0475 \\ -0.0036 & 0.455 & -0.125 & -0.265 \\ -8.18 & 23.0 & -159 & 49.4 \\ 0.123 & -0.345 & 2.19 & -2.24 \end{pmatrix}. \quad (37)$$

the term relevant to D meson mixing is very small

$$(V_L \mathcal{M}_u'' V_R^\dagger)_{21} \simeq -3.6 \text{ MeV}. \quad (38)$$

Then the effective Lagrangian reads,

$$\mathcal{L}_{\text{eff}} \sim \frac{G_F}{m_{H_1}^2} (V_L \mathcal{M}_u'' V_R^\dagger)_{21}^2 \left[(c\bar{u})^2 - (c\gamma_5 \bar{u})^2 \right], \quad (39)$$

which produces a extremely small contribution to $D^0 - \bar{D}^0$ mixing,

$$\begin{aligned} \Delta m_D &\sim \langle \bar{D} | \mathcal{L}_{\text{eff}} | D \rangle \\ &\sim \frac{G_F}{m_{H_1}^2} (V_L \mathcal{M}_u'' V_R^\dagger)_{21}^2 m_D F_D^2 \left(\frac{m_D}{m_c} \right)^2 \\ &\sim 2.5 \times 10^{-17} \left(\frac{\text{TeV}}{m_{H_1}} \right)^2 \text{ GeV}. \end{aligned} \quad (40)$$

The experimental value of the mass difference is of order 10^{-14} GeV; this therefore does not severely constrain the H_1 mass. The right panel of Fig. 2 shows that the flavor-changing contribution to $D^0 - \bar{D}^0$ mixing is stable against the variation of the parameters in the potential and variation of v_R . One might naïvely expect that $(V_L \mathcal{M}_u'' V_R^\dagger)_{21}$ should be dominated by the term $\sim m_c (V_R^\dagger)_{21}$, however, numerical data shows that this term is always cancelled by other contributions, e.g. the terms proportional to $(\mathcal{M}_u'')_{33}$ or $(\mathcal{M}_u'')_{34}$, leaving a very tiny contribution. As stated in the introduction, this is guaranteed by the fact that the SM quark mixing is produced by tiny mixing with the vector-like heavy fermion.

V. LHC PROSPECTS FOR H_1

In our model, H_1 behave somewhat like a heavy copy of the SM Higgs: it has hierarchical couplings to the SM fermions, and couples also to the gauge bosons. Numerical analysis

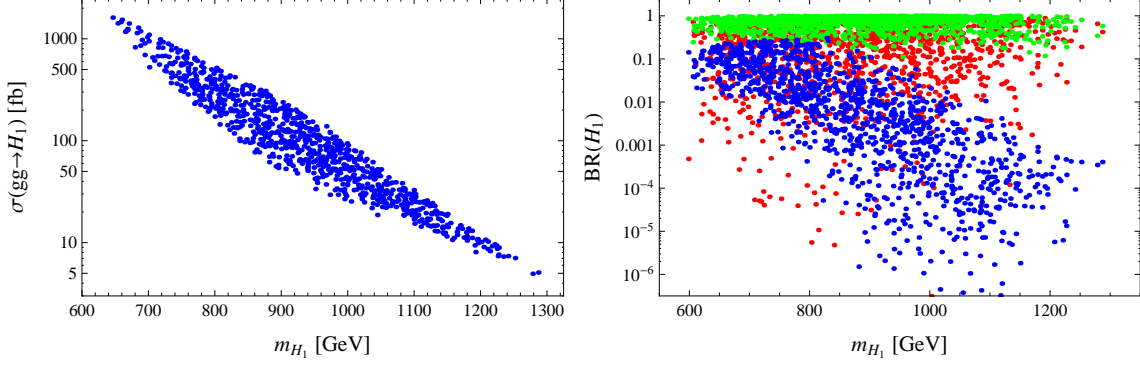


FIG. 4: Left panel: H_1 production cross section $\sigma(gg \rightarrow H_1)$ at LHC with center-of-mass energy of 14 TeV, as function of H_1 mass. Right panel: Branching ratios of H_1 decay, as function of its mass. The decay channels $H_1 \rightarrow hh$, $t\bar{t}$, WW are shown, respectively, as red, green and blue spots.

reveals that its coupling to the top quark is generally of order one; therefore, similarly to the SM Higgs, the top-loop induced gluon fusion process would be the dominant production channel for H_1 at the LHC. It is very straightforward to obtain the leading order cross section for the production process $gg \rightarrow H_1$: we need only to rescale that for the SM Higgs (extrapolated to \sim TeV mass) by the Yukawa coupling ratio $(V_L \mathcal{M}_u'' V_R^\dagger)_{33}^2 / (V_L \mathcal{M}_u' V_R^\dagger)_{33}^2$. The left panel of Fig. 4 depicts the sizable production cross section and its dependence on the scalar mass. It is evident in the figure that, for a TeV H_1 , the cross section can be as large as 100 fb at $\sqrt{s} = 14$ TeV, and we could expect thousands of H_1 events produced at upgraded LHC. For a lighter H_1 , the cross section could be even larger.

We now sketch how the second Higgs decays after its production. Once the scalars obtain their non-zero vevs, there appear cubic couplings among the Higgs states, which in the flavor basis are given by,

$$\begin{aligned}
& \left[\alpha'_1 \kappa_1 (\text{Re}\phi_1^0)^3 + \alpha'_1 \kappa_2 (\text{Re}\phi_2^0)^3 + \beta' v_\rho (\text{Re}\rho^0)^3 + \gamma' v_R (\text{Re}\delta_R^0)^3 \right] \\
& + \frac{1}{2} \left[x'_1 v_\rho (\text{Re}\phi_1^0)^2 (\text{Re}\rho^0) + x'_1 \kappa_1 (\text{Re}\phi_1^0) (\text{Re}\rho^0)^2 + x'_2 v_\rho (\text{Re}\phi_2^0)^2 (\text{Re}\rho^0) + x'_2 \kappa_2 (\text{Re}\phi_2^0) (\text{Re}\rho^0)^2 \right] \\
& + \frac{1}{2} \left[y'_1 v_R (\text{Re}\phi_1^0)^2 (\text{Re}\delta_R^0) + y'_1 \kappa_1 (\text{Re}\phi_1^0) (\text{Re}\delta_R^0)^2 + y'_2 v_R (\text{Re}\phi_2^0)^2 (\text{Re}\delta_R^0) + y'_2 \kappa_2 (\text{Re}\phi_2^0) (\text{Re}\delta_R^0)^2 \right] \\
& + \frac{1}{2} \left[z' v_R (\text{Re}\rho^0)^2 (\text{Re}\delta_R^0) + z' v_\rho (\text{Re}\rho^0) (\text{Re}\delta_R^0)^2 \right] \\
& - \frac{1}{\sqrt{2}} \left[M' (\text{Re}\phi_1^0) (\text{Re}\rho^0) (\text{Re}\delta_R^0) + M' (\text{Re}\phi_2^0) (\text{Re}\rho^0) (\text{Re}\delta_R^0) \right].
\end{aligned} \tag{41}$$

After rotating these states to the physical basis, we can easily get the dimensionful cubic coupling $m_{H_1 hh} H_1 hh$ from the potential. To calculate the decay width $\Gamma(H_1 \rightarrow hh)$, we assume all the relevant original couplings in the potential (those which are not constrained by the observation that $m_h = 125$ GeV) lie in the perturbative range⁴ $\sim (0, 3)$. This leads

⁴ If these parameters have much larger values, e.g. $\lesssim 4\pi$, then in a large portion of parameter space, the decay channel $H_1 \rightarrow hh$ dominates over others. However, here we consider the preferable smaller values of couplings, $\sim (0, 3)$ in giving the branching ratios. In numerical calculations we take values in the exact range $[0, \sqrt{4\pi}]$.

to the decay width as:

$$\Gamma(H_1 \rightarrow hh) = \frac{1}{8\pi} \frac{m_{H_1 hh}^2}{m_{H_1}} \left(1 - \frac{4m_h^2}{m_{H_1}^2}\right)^{1/2}. \quad (42)$$

Turning to the fermion decay channels, the top quark mode dominates over others. As for the coupling of H_1 to top quark, it is easy to see that the Yukawa coupling $y_{H_1 t\bar{t}} = \sqrt{2}(V_L M_u'' V_R^\dagger)_{33}/v$, which produces the fermionic decay width

$$\Gamma(H_1 \rightarrow t\bar{t}) = \frac{3}{16\pi} \cdot |y_{H_1 t\bar{t}}|^2 m_{H_1} \left(1 - \frac{4m_t^2}{m_{H_1}^2}\right)^{3/2}. \quad (43)$$

Besides the scalars and fermions H_1 can also decay into SM gauge bosons, WW , ZZ , but those widths are generally suppressed by the smaller gauge coupling. For instance, for $H_1 \rightarrow WW$, we define $m_{H_1 WW} = U_{21}\kappa_1 + U_{22}\kappa_2 + U_{23}v_\rho$ and $f_W = \frac{1}{2}g^2 m_{H_1 WW}$, then the width reads

$$\Gamma(H_1 \rightarrow WW) = \frac{1}{8\pi} \frac{f_W^2}{m_{H_1}} \left[1 + \frac{1}{2} \left(1 - \frac{m_{H_1}^2}{2m_W^2}\right)^2\right] \left(1 - \frac{4m_W^2}{m_{H_1}^2}\right)^{1/2}. \quad (44)$$

The ZZ channel is expected to have similar width. As an explicit numerical example, the branching ratios of these different decay channels are shown in the right panel of Fig. 4. Obviously, in almost all the parameter space, the top quark channel dominates as mentioned before, enhanced by the order one Yukawa coupling.

VI. LEPTON SECTOR

Let us briefly comment on the lepton sector of the model. It is clear from Eq. (5) that $M_\ell = rM_D$, where $M_{\ell,D}$ are the charged lepton and neutrino Dirac masses. However, due to this fact and the fact that $v_R \sim 3$ TeV, simple generic type I seesaw cannot reproduce the observed neutrino spectrum and mixings. We have tried possible new textures for y_R matrix but have not found any that will help us to get right neutrino mass pattern. It seems that the simplest way to accommodate small neutrino masses in our scheme is to add three gauge singlet neutrinos, a right handed doublet Higgs field and invoke inverse seesaw mechanism [13]. According to the inverse seesaw paradigm, the right handed neutrino of the left-right model is a quasi-Dirac neutrino and its characteristic signature is a trilepton final state rather than the $\ell^\pm \ell^\pm jj$ mode. This implies that the search strategy for the W_R at LHC must change in case a TeV scale heavy Higgs is found. This situation can also have interesting implications for neutrinoless double beta decay [24].

VII. OTHER PHENOMENOLOGICAL AND THEORETICAL COMMENTS

(i) In addition to the three near TeV LHC accessible neutral Higgs fields, two new fermions are the top partner t' and the 5/3 charged quarks with new phenomenology. The presence of the 5/3 charged quark also alters the phenomenology of the doubly charged Higgs bosons [25]. From our analysis, it appears that the CKM fit requires $m_{t'} \sim 3$ TeV. Since we expect

$m_{t'} \simeq m_{Q_{5/3}}$ due to the vector-like nature of the extra fermion doublet, both the t' and $Q_{5/3}$ could be accessible at the LHC. As was noted in [25], unlike the decay properties of $Q_{5/3}$ discussed in the literature so far [26], where it is assumed that $Q_{5/3} \rightarrow t + W^+$, in our scheme (as in the scheme in [25]), the primary decay mode of $Q_{5/3}$ could be $Q_{5/3} \rightarrow \Delta^{++} + d$ and/or $Q_{5/3} \rightarrow \rho^{++} + d$ leading to different LHC signal.

(ii) The model has also new charged scalar states beyond that of the minimal LRSM. For instance there are the new doubly charged fields ρ^{++} , whose mass is expected to be in the TeV range. Its likely decay mode is $\rho^{++} \rightarrow \Delta^+ \phi_2^+$ with Δ^+ subsequently decaying to leptonic final states. There are also new singly charged states as partners of $H_{1,2}$ with associated rich phenomenology as are the pseudo-scalar partners of $H_{1,2}$. We do not pursue this any further in this paper.

(iii) An interesting point about our model is that even though the SM Higgs coupling to $b\bar{b}$ arises in an indirect way compared to SM (via coupling to ϕ_2^0 and fraction of SM Higgs in ϕ_2^0), its magnitude remains almost same as in SM.

(iv) It is also worth noting that one could envision adding more than one vector-like quark multiplets (instead of one that we have added). Such models would lead to changes in the details of the model e.g. Higgs decays, masses etc. The model discussed here should be considered as an existence of proof of models with lower Higgs mass in left-right symmetric models rather than as a definite final model.

(v) Finally, a theoretical comment on the model: the new fermion and scalar multiplets we have chosen to include can emerge from an $SO(10)$ grand unified theory. For instance, the fermion Q' and Higgs fields ρ added to minimal LRSM could also have grand unified origin since the Q' is part of the **560**+**560**^{*} representation whereas the ρ -Higgs field is a sub-multiplet of **210** representation. The new Yukawa couplings involving Q' and ρ in Eq. (5), y_f and y_g couplings can arise from $SO(10)$ invariant couplings **16** · **210** · **560**^{*} and **16** · **126**^{*} · **560** in a possible GUT version of the theory. We have also checked that if there is a complex $Y = 0$ $SU(2)_L$ triplet in the theory at the TeV scale, three couplings of SM almost unify at the one loop level via the chain $G_{SM} \rightarrow G_{LRSM} \rightarrow SO(10)$ with G_{LRSM} at the TeV scale. Given that we have not included threshold effects as well as two loop contributions, this unification is likely to be better. Such triplets could be the ones present in **45** or **54** Higgs fields that are sued to break the $SO(10)$ symmetry. We assume that all the other fields in the fermion multiplet **560** except the Q' must be at the GUT scale. We do not pursue these and other detailed aspects of grand unification possibility any further since it is beyond the scope of this paper.

VIII. SUMMARY

In this paper, we have explored the question of the second Higgs mass in the left-right symmetric models. The well known fact that in the minimal version of this model, the second neutral Higgs mass is more than 12-15 TeV implies that discovery of a neutral Higgs with a few TeV mass ($\ll 10$ TeV) would rule out the minimal LRMS. The question would then be: should the search for W_R at LHC continue in this case? In other words, are there versions of the left-right model that keep the W_R mass in the 5-6 TeV range while at the same time having a Higgs mass of a few TeV or less without conflicting with meson-anti-meson constraints ($K^0 - \bar{K}^0$, $B^0 - \bar{B}^0$ mixing etc)? We have provided an existence proof

of such models in this paper and what their basic features should be. It appears from our example that such models should have heavier vector like quarks and/or extra Higgs fields to generate desired CKM mixings. An interesting consequence of these models also appears to be that the search mode for W_R at the LHC should change to $\ell^\pm \ell^\mp \ell^\pm + \text{missing E}$ rather than $\ell^\pm \ell^\pm jj$ with no missing E, which is currently being pursued [10]. In the particular example we provide, the heavier neutral Higgs masses are all below the W_R mass, although their precise mass values could shift in this range depending on model details. Needless to say that search for a TeV scale Higgs is therefore crucial for understanding the left-right symmetric extensions of the standard model and how neutrino masses arise in such models.

Acknowledgement

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Appendix A: CKM fit without CP violation

In this appendix, as a toy test of our model, we consider the fit of CKM matrix without any CP violation. Counting the numbers of parameters in Eq. (11): on the LHS, we have eight parameters r , $f_{1,2,3}$, $g_{1,2,3}$ and M , whereas on the RHS, we have ten to be fit, the four quark masses $m_{u,c,t}$, m'_t and the six independent mixing angles $\theta_{12,23,13,14,24,34}$. Of the ten parameters on the RHS, six are known and the other four (m'_t and $\theta_{14,24,34}$) involve new physics are yet undetermined. To find reasonable solutions for our model, we need to fix two of the four unknown parameters on the RHS. Explicitly, we choose⁵

$$\begin{aligned} m'_t &= 3 \text{ TeV} , \\ \theta_{34} &= 0.016 . \end{aligned} \tag{A1}$$

The choice of the value for θ_{34} means that it lies on the 2σ boundary of the SM constraint.

A typical solution for the parameters that fits all quark masses and mixings is given below (all masses in unit of GeV):

$$\begin{aligned} r &= 10.1 , \quad f_1 = -0.916 , \quad g_1 = -465 , \quad M = 668 . \\ f_2 &= 6.26 , \quad g_2 = -2260 , \\ f_3 &= -157 , \quad g_3 = -1797 , \end{aligned} \tag{A2}$$

⁵ We find that if we set m'_t at lower values, say 1 TeV, then, due to mixing of t' with the SM fermions, the values of M would be much smaller than the TeV scale (e.g., for $m'_t = 1 \text{ TeV}$, $M \sim 10 \text{ GeV}$), which seems less appealing. Thus we choose a somewhat larger t' quark mass.

The resulting 4×4 quark mixing matrix reads

$$\begin{pmatrix} 0.974 & 0.225 & 0.0035 & 0.000047 \\ -0.225 & 0.973 & 0.041 & 0.00032 \\ 0.0058 & -0.0409 & 0.999 & 0.016 \\ -0.000069 & 0.00034 & -0.0165 & 0.999 \end{pmatrix}, \quad (\text{A3})$$

which fits very well the 3×3 CKM matrix and it is transparent that the mixings between the SM up-type quark with the 4th one are very small.

$$\begin{aligned} s_{14} &= 0.000047, \\ s_{24} &= 0.00032, \\ s_{34} &= 0.016, \end{aligned} \quad (\text{A4})$$

and it is reasonable that the mixing with the top quark is comparatively much larger than with the first two generations.

Appendix B: Potential and scalar mass-squared matrices

1. Potential and minimization conditions

In this appendix, we discuss the Higgs potential of our model, its minimization to get the mass eigenstates and eigenvalues for the CP even and CP odd real scalar fields of the model. The full Higgs potential of the model invariant under the discrete and gauge symmetries is:

$$\begin{aligned} V = & -\mu_1^2 \text{Tr}(\phi^\dagger \phi) + M_\rho^2 \text{Tr}(\rho^\dagger \rho) - \mu_3^2 \text{Tr}(\Delta_R \Delta_R^\dagger) \\ & + \alpha_1 [\text{Tr}(\phi^\dagger \phi)]^2 + \alpha_2 \left\{ [\text{Tr}(\tilde{\phi} \phi^\dagger)]^2 + [\text{Tr}(\tilde{\phi}^\dagger \phi)]^2 \right\} + \alpha_3 \text{Tr}(\tilde{\phi} \phi^\dagger) \text{Tr}(\tilde{\phi}^\dagger \phi) \\ & + \alpha_4 \text{Tr}(\phi^\dagger \phi \phi^\dagger \phi) + \alpha_5 \text{Tr}(\phi^\dagger \phi \tilde{\phi}^\dagger \tilde{\phi}) + \alpha_6 [\text{Tr}(\phi^\dagger \tilde{\phi} \phi^\dagger \tilde{\phi}) + \text{Tr}(\phi \tilde{\phi}^\dagger \phi \tilde{\phi}^\dagger)] \\ & + \beta_1 [\text{Tr}(\rho^\dagger \rho)]^2 + \beta_2 \text{Tr}(\tilde{\rho} \rho^\dagger) \text{Tr}(\tilde{\rho}^\dagger \rho) + \beta_3 \text{Tr}(\rho^\dagger \rho \rho^\dagger \rho) + \beta_4 \text{Tr}(\rho^\dagger \rho \tilde{\rho}^\dagger \tilde{\rho}) \\ & + \gamma_1 [\text{Tr}(\Delta_R \Delta_R^\dagger)]^2 + \gamma_2 \text{Tr}(\Delta_R \Delta_R) \text{Tr}(\Delta_R^\dagger \Delta_R^\dagger) \\ & + \gamma_3 \text{Tr}(\Delta_R \Delta_R^\dagger \Delta_R \Delta_R^\dagger) + \gamma_4 \text{Tr}(\Delta_R \Delta_R \Delta_R^\dagger \Delta_R^\dagger) \\ & + x_1 \text{Tr}(\phi^\dagger \phi) \text{Tr}(\rho^\dagger \rho) + x_2 \text{Tr}(\phi^\dagger \phi \rho^\dagger \rho) + x_3 \text{Tr}(\phi \phi^\dagger \rho \rho^\dagger) \\ & + x_4 \text{Tr}(\tilde{\phi} \tilde{\phi}^\dagger \rho \rho^\dagger) + x_5 \text{Tr}(\tilde{\phi}^\dagger \tilde{\phi} \rho^\dagger \rho) \\ & + x_6 [\text{Tr}(\tilde{\phi} \tilde{\rho}^\dagger \phi \rho^\dagger) + \text{Tr}(\tilde{\rho} \tilde{\phi}^\dagger \rho \phi^\dagger)] + x_7 [\text{Tr}(\phi \tilde{\rho}^\dagger \tilde{\phi} \rho^\dagger) + \text{Tr}(\rho \tilde{\phi}^\dagger \tilde{\rho} \phi^\dagger)] \\ & + y_1 \text{Tr}(\phi^\dagger \phi) \text{Tr}(\Delta_R \Delta_R^\dagger) + y_2 \text{Tr}(\phi^\dagger \phi \Delta_R^\dagger \Delta_R) + y_3 \text{Tr}(\phi^\dagger \phi \Delta_R \Delta_R^\dagger) \\ & + z_1 \text{Tr}(\rho^\dagger \rho) \text{Tr}(\Delta_R \Delta_R^\dagger) + z_2 \text{Tr}(\rho^\dagger \rho \Delta_R^\dagger \Delta_R) + z_3 \text{Tr}(\rho^\dagger \rho \Delta_R \Delta_R^\dagger). \end{aligned} \quad (\text{B1})$$

To get non-vanishing vev of the extra ρ scalar, we introduce the terms soft breaking the discrete symmetry,

$$M' [\text{Tr}(\tilde{\rho} \Delta_R \phi^\dagger) + \text{Tr}(\tilde{\rho} \Delta_R \tilde{\phi}^\dagger) + \text{h.c.}]. \quad (\text{B2})$$

For simplicity, we assume all the couplings in the Lagrangian are real parameters.

The four minimization conditions are:

$$\frac{\partial}{\partial \kappa_1} V = \frac{\partial}{\partial \kappa_2} V = \frac{\partial}{\partial v_\rho} V = \frac{\partial}{\partial v_R} V = 0. \quad (\text{B3})$$

They lead to the following relations among the vevs and the coefficient in the potential,

$$\begin{aligned} \frac{\mu_1^2}{v_R^2} &= \frac{y'_1}{2} + \alpha'_1 \frac{\kappa_1^2}{v_R^2} + \alpha'_2 \frac{\kappa_2^2}{v_R^2} + \frac{x'_1}{2} \frac{v_\rho^2}{v_R^2} - \frac{1}{\sqrt{2}} \frac{M'}{\kappa_1} \frac{v_\rho}{v_R}, \\ \frac{\mu_2^2}{v_R^2} &= \frac{y'_2}{2} + \alpha'_2 \frac{\kappa_1^2}{v_R^2} + \alpha'_1 \frac{\kappa_2^2}{v_R^2} + \frac{x'_2}{2} \frac{v_\rho^2}{v_R^2} - \frac{1}{\sqrt{2}} \frac{M'}{\kappa_2} \frac{v_\rho}{v_R}, \\ \frac{1}{\sqrt{2}} \frac{\kappa_1 + \kappa_2}{v_R} \frac{M'}{v_\rho} - \frac{M_\rho^2}{v_R^2} &= \frac{z'}{2} + \frac{x'_1}{2} \frac{\kappa_1^2}{v_R^2} + \frac{x'_2}{2} \frac{\kappa_2^2}{v_R^2} + \beta' \frac{v_\rho^2}{v_R^2}, \\ \frac{\mu_3^2}{v_R^2} &= \gamma' + \frac{y'_1}{2} \frac{\kappa_1^2}{v_R^2} + \frac{y'_2}{2} \frac{\kappa_2^2}{v_R^2} + \frac{z'}{2} \frac{v_\rho^2}{v_R^2} - \frac{1}{\sqrt{2}} \frac{M'}{v_R} \frac{\kappa_1 + \kappa_2}{v_R} \frac{v_\rho}{v_R}, \end{aligned} \quad (\text{B4})$$

where

$$\begin{aligned} \alpha'_1 &= \alpha_1 + \alpha_4, & x'_1 &= x_1 + x_2 + x_4 - 2x_6, \\ \alpha'_2 &= \alpha_1 + 4\alpha_2 + 2\alpha_3 + \alpha_5 + 2\alpha_6, & x'_2 &= x_1 + x_3 + x_5 - 2x_7, \\ \beta' &= \beta_1 + \beta_3, & y'_1 &= y_1 + y_2, \\ \gamma' &= \gamma_1 + \gamma_3, & y'_2 &= y_1 + y_3, \\ & & z' &= z_1 + z_2. \end{aligned} \quad (\text{B5})$$

Assuming the two mass parameters $M_\rho, M' \sim v_{EW}$, we find at leading order of $\kappa_1, \kappa_2, v_\rho \ll v_R$ that these relations are greatly simplified,

$$\begin{aligned} \frac{\mu_1^2}{v_R^2} &\simeq \frac{y_1 + y_2}{2}, \\ \frac{\mu_2^2}{v_R^2} &\simeq \frac{y_1 + y_3}{2}, \\ \frac{1}{\sqrt{2}} \frac{\kappa_1 + \kappa_2}{v_R} \frac{M'}{v_\rho} &\simeq \frac{z_1 + z_2}{2}, \\ \frac{\mu_3^2}{v_R^2} &\simeq \gamma_1 + \gamma_3. \end{aligned} \quad (\text{B6})$$

The first two equations imply (at the leading order of $\frac{\kappa_{1,2}, v_\rho}{v_R}$) $y_2 \simeq y_3$. There are two equations (the first and fourth ones) for the vev v_R , indicating fine-tuning of some of the coefficients, as in the case of manifest minimal LRSM [7].

2. Neutral Higgs boson mass-squared matrices

Using the above conditions, we find that in the basis of $\{\text{Re}\phi_1^0, \text{Re}\phi_2^0, \text{Re}\rho^0, \text{Re}\delta_R^0\}$, the neutral Higgs mass-square matrix elements read [17]

$$M_{11}^{2\text{Re}} = -\mu_1^2 + 3\alpha'_1 \kappa_1^2 + \alpha'_2 \kappa_2^2 + \frac{x'_1}{2} v_\rho^2 + \frac{y'_1}{2} v_R^2,$$

$$\begin{aligned}
M_{22}^{2\text{Re}} &= -\mu_1^2 + \alpha'_2 \kappa_1^2 + 3\alpha'_1 \kappa_2^2 + \frac{x'_2}{2} v_\rho^2 + \frac{y'_2}{2} v_R^2, \\
M_{33}^{2\text{Re}} &= +M_\rho^2 + \frac{x'_1}{2} \kappa_1^2 + \frac{x'_2}{2} \kappa_2^2 + 3\beta' v_\rho^2 + \frac{z'}{2} v_R^2, \\
M_{44}^{2\text{Re}} &= -\mu_3^2 + \frac{y'_1}{2} \kappa_1^2 + \frac{y'_2}{2} \kappa_2^2 + \frac{z'}{2} v_\rho^2 + 3\gamma' v_R^2, \\
M_{12}^{2\text{Re}} &= 2\alpha'_2 \kappa_1 \kappa_2, \\
M_{13}^{2\text{Re}} &= x'_1 \kappa_1 v_\rho - \frac{1}{\sqrt{2}} M' v_R, \\
M_{23}^{2\text{Re}} &= x'_2 \kappa_2 v_\rho - \frac{1}{\sqrt{2}} M' v_R, \\
M_{14}^{2\text{Re}} &= y'_1 \kappa_1 v_R - \frac{1}{\sqrt{2}} M' v_\rho, \\
M_{24}^{2\text{Re}} &= y'_2 \kappa_2 v_R - \frac{1}{\sqrt{2}} M' v_\rho, \\
M_{34}^{2\text{Re}} &= z' v_\rho v_R - \frac{1}{\sqrt{2}} M' (\kappa_1 + \kappa_2), \tag{B7}
\end{aligned}$$

while in the basis of $\{\text{Im}\phi_1^0, \text{Im}\phi_2^0, \text{Im}\rho^0, \text{Im}\delta_R^0\}$, the elements for the pseudoscalar mass-square matrix are

$$\begin{aligned}
M_{11}^{2\text{Im}} &= -\mu_1^2 + \alpha'_1 \kappa_1^2 + \alpha''_2 \kappa_2^2 + \frac{x'_1}{2} v_\rho^2 + \frac{y'_1}{2} v_R^2, \\
M_{22}^{2\text{Im}} &= -\mu_1^2 + \alpha''_2 \kappa_1^2 + \alpha'_1 \kappa_2^2 + \frac{x'_2}{2} v_\rho^2 + \frac{y'_2}{2} v_R^2, \\
M_{33}^{2\text{Im}} &= +M_\rho^2 + \frac{x'_1}{2} \kappa_1^2 + \frac{x'_2}{2} \kappa_2^2 + \beta' v_\rho^2 + \frac{z'}{2} v_R^2, \\
M_{44}^{2\text{Im}} &= -\mu_3^2 + \frac{y'_1}{2} \kappa_1^2 + \frac{y'_2}{2} \kappa_2^2 + \frac{z'}{2} v_\rho^2 + \gamma' v_R^2, \\
M_{12}^{2\text{Im}} &= -4\alpha''_1 \kappa_1 \kappa_2, \\
M_{13}^{2\text{Im}} &= +\frac{1}{\sqrt{2}} M' v_R, \\
M_{23}^{2\text{Im}} &= -\frac{1}{\sqrt{2}} M' v_R, \\
M_{14}^{2\text{Im}} &= -\frac{1}{\sqrt{2}} M' v_\rho, \\
M_{24}^{2\text{Im}} &= +\frac{1}{\sqrt{2}} M' v_\rho, \\
M_{34}^{2\text{Im}} &= -\frac{1}{\sqrt{2}} M' (\kappa_1 + \kappa_2), \tag{B8}
\end{aligned}$$

with

$$\begin{aligned}
\alpha''_1 &= 2\alpha_2 + \alpha_6, \\
\alpha''_2 &= \alpha_1 - 4\alpha_2 + 2\alpha_3 + \alpha_5 - 2\alpha_6. \tag{B9}
\end{aligned}$$

When the minimization conditions are applied, the mass-square matrix for the pseudoscalars is greatly simplified,

$$\begin{pmatrix} -4\alpha_1''\kappa_2^2 + \frac{M'v_\rho v_R}{\sqrt{2}\kappa_1} & -4\alpha_1''\kappa_1\kappa_2 & \frac{M'v_R}{\sqrt{2}} & -\frac{M'v_\rho}{\sqrt{2}} \\ -4\alpha_1''\kappa_1\kappa_2 & -4\alpha_1''\kappa_1^2 + \frac{M'v_\rho v_R}{\sqrt{2}\kappa_2} & -\frac{M'v_R}{\sqrt{2}} & \frac{M'v_\rho}{\sqrt{2}} \\ \frac{M'v_R}{\sqrt{2}} & -\frac{M'v_\rho}{\sqrt{2}} & \frac{M'(\kappa_1+\kappa_2)v_R}{\sqrt{2}v_\rho} & -\frac{M'(\kappa_1+\kappa_2)}{\sqrt{2}} \\ -\frac{M'v_\rho}{\sqrt{2}} & \frac{M'v_\rho}{\sqrt{2}} & -\frac{M'(\kappa_1+\kappa_2)}{\sqrt{2}} & \frac{M'(\kappa_1+\kappa_2)v_\rho}{\sqrt{2}v_R} \end{pmatrix}. \quad (\text{B10})$$

As we expect, This matrix has two massless Goldstone boson states, one of which is predominately from $(\text{Im}\delta_R^0)$ and is “eaten” by Z' and the other one from combination of the imaginary parts of $\phi_{1,2}^0$ and ρ^0 becoming the longitudinal component of Z . The lighter one of the two massive states, namely A_1 , is expected to lie at the TeV scale, and the other one A_2 being a bit heavier, as explicitly shown in Fig. 1.

At the leading order in the parameters $\frac{\kappa_1, \kappa_2, v_\rho}{v_R}$, with help of the equations (B6), the mass-square matrix for the CP-even scalars comes out very simple,

$$M^{2\text{Re}} = \text{diag}\{0, 0, 0, 2\gamma'v_R^2\}, \quad (\text{B11})$$

implying that of the four CP-even scalars, only one is at the right-handed scale, and the other three are at lower scales. When we include the next-to-leading order of $\frac{\kappa_1, \kappa_2, v_\rho}{v_R}$ terms, the matrix becomes

$$M^{2\text{Re}} = v_R^2 \begin{pmatrix} \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_1 v_R} & 0 & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & y_1' \frac{\kappa_1}{v_R} \\ 0 & \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_2 v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & y_2' \frac{\kappa_2}{v_R} \\ -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & \frac{1}{\sqrt{2}} \frac{\kappa_1+\kappa_2}{v_R} & z' \frac{v_\rho}{v_R} \\ y_1' \frac{\kappa_1}{v_R} & y_2' \frac{\kappa_2}{v_R} & z' \frac{v_\rho}{v_R} & 2\gamma' \end{pmatrix}, \quad (\text{B12})$$

in which $y_1'\kappa_1 = y_2'\kappa_2$, and we have set $M' = v_\rho$ for simplicity⁶. As the right-handed scale is much higher than the electroweak scale $\kappa_1, \kappa_2, v_\rho$, the masses for three lighter scalars are mainly from the upper-left 3×3 block of the matrix (corresponding to the scalars $\text{Re}\phi_{1,2}^0$ and $\text{Re}\phi^0$),

$$v_R^2 \begin{pmatrix} \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_1 v_R} & 0 & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} \\ 0 & \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_2 v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} \\ -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & \frac{1}{\sqrt{2}} \frac{\kappa_1+\kappa_2}{v_R} \end{pmatrix}. \quad (\text{B13})$$

One can easily find that one of the eigenstates is massless in the direction of $(\kappa_1, \kappa_2, v_\rho)$, then there is always a state in the full 4×4 matrix with negative values of mass square (this mass eigenstate is expected to play the role of the “light” SM Higgs with mass of 125 GeV, after higher order terms are included that makes this negative value positive). The two remaining states are expected to be at the scale of $\sqrt{vv_R}$. To get the four mass eigenstates all with positive masses-square values (especially the 125 GeV SM Higgs), we

⁶ There are enough parameters in the potential to make such a choice possible.

need to expand the mass-square matrix to the next order,

$$M^{2\text{Re}} = v_R^2 \begin{pmatrix} \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_1 v_R} & \frac{2\alpha'_2 \kappa_1 \kappa_2}{v_R^2} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & y'_1 \frac{\kappa_1}{v_R} \\ \frac{2\alpha'_2 \kappa_1 \kappa_2}{v_R^2} & \frac{1}{\sqrt{2}} \frac{v_\rho^2}{\kappa_2 v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & y'_2 \frac{\kappa_2}{v_R} \\ -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & -\frac{1}{\sqrt{2}} \frac{v_\rho}{v_R} & \frac{1}{\sqrt{2}} \frac{\kappa_1 + \kappa_2}{v_R} & z' \frac{v_\rho}{v_R} \\ y'_1 \frac{\kappa_1}{v_R} & y'_2 \frac{\kappa_2}{v_R} & z' \frac{v_\rho}{v_R} & 2\gamma' \end{pmatrix}. \quad (\text{B14})$$

The new terms in the (1, 2) (and (2, 1)) element would compensate the negative contribution from mixing with fourth heavy state and then produce a positive value for mass-square of the would-be lightest state (the SM Higgs), thus the quartic coupling combination α'_2 is expected to have large value.

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